

Symmetry Principles toward Solutions of the μ Problem

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Abstract

We stress that a natural solution of the μ problem requires two ingredients: a symmetry that would enforce $\mu = 0$ as well as the occurrence of a small breaking parameter that generates a nonzero μ . It is suggested that both the Peccei-Quinn symmetry and the spontaneously broken R symmetry may be the sources of the needed μ term in the minimal supersymmetric standard model provided that they are spontaneously broken at a scale $10^{10} - 10^{12}$ GeV. To solve the strong CP problem with a hidden sector confining group, both of these symmetries are needed in superstring models with an anomalous $U(1)_A$.

1. Introduction

It is often assumed that one can understand many aspects of the low energy electroweak phenomena from supergravity interactions at high energies governed by an underlying superstring theory [1]. Even though there does not exist a *standard* superstring model, it is a general expectation that one of numerous string vacua might coincide with the minimal supersymmetric standard model (MSSM) in supergravity [2]. Nevertheless, the attractive MSSM has a few outstanding theoretical parameter problems: the cosmological constant problem [3], the strong CP problem [4], and the μ problem [5]. From the top-down approach, these may be solved by the *theory*. But from the bottom-up approach, these need explanations and constitute very interesting and challenging problems. In this paper, we address two of these problems: the strong CP problem and the μ problem.

The strong CP problem seems to have an inherent solution in string models due to the existence of the model-independent axion [6]. However, the model-independent axion is known to have the axion decay constant problem if it is designed to solve the strong CP problem [7]. Apart from this cosmological problem, the need for a hidden sector confining group to break supersymmetry at the intermediate scale requires another axion to settle the QCD vacuum angle at zero. However, the string models are known to have no continuous global symmetry. Therefore, there does not exist any room for the additional axion at the string level. Thus the prospect for a solution of the strong CP problem through an invisible axion looks doomed.

On the other hand, the MSSM admits

$$W_\mu = \mu H_1 H_2 \tag{1}$$

in the superpotential, where H_1 and H_2 are two Higgs doublets with $Y = -1/2$ and $Y = 1/2$ and μ is a free parameter whose magnitude is in principle different from the magnitude of $m_{3/2} \sim M_{SUSY}$ characterizing the soft SUSY breaking terms in supergravity. Electroweak phenomenology suggests that μ is nonzero and falls in the 100 GeV region similar to the magnitude of soft parameters. In addition, axion phenomenology suggests that μ cannot be

zero, or there results an unwanted 0.1 MeV axion. Theoretically, however, one expects that μ is of order the Planck scale unless it is forbidden by some symmetry. At present, there exist several suggestions toward a solution of the μ problem [5,8,9].

In this paper, we will try to understand these two outstanding problems from symmetry principles, which will work as a guideline for future construction of superstring models.

2. The Peccei–Quinn and R Symmetries

From the early days of the μ problem [5], it was suggested that a symmetry, presumably a Peccei–Quinn symmetry with an invisible axion [10], may be needed toward a solution of the problem. The reason is simple. Consider the following Yukawa super-potential

$$W_Y = d_L^{cT} F_d q_L H_1 + u_L^{cT} F_u q_L H_2 \quad (2)$$

where $F_{d,u}$ is the 3×3 Yukawa coupling matrix, q_L is a column vector of three SU(2) doublet superfields, d_L^c and u_L^c are column vectors of up- and down-type anti-quark superfields. A Peccei–Quinn symmetry¹

$$H_1 \longrightarrow e^{i\alpha} H_1, \quad H_2 \longrightarrow e^{i\alpha} H_2, \quad \{q_L, d_L^c, u_L^c\} \longrightarrow e^{-\frac{i}{2}\alpha} \{q_L, d_L^c, u_L^c\} \\ W \longrightarrow W. \quad (3)$$

forbids W_μ . One can introduce the Peccei–Quinn symmetry to understand the smallness of μ , and break it at $10^{10} - 10^{12}$ GeV, generating an electroweak scale μ through nonrenormalizable interactions generated by gravity [5]. The Peccei–Quinn symmetry seems to be a necessity to fix the scale of μ at M_{SUSY} .

In this paper, we also suggest that R symmetry can be used as another symmetry in addition to the Peccei–Quinn symmetry.² With $\mu = 0$, we assign an R symmetry,

$$\{H_1, H_2\} \longrightarrow \{H_1, H_2\},$$

¹For the heavy quark axion, $H_{1,2} \rightarrow H_{1,2}$, and hence it cannot be used for the small μ term.

²The use of R symmetry for the μ term has been suggested before [11,12].

$$\begin{aligned}\{q_L, d_L^c, u_L^c\} &\longrightarrow e^{i\beta}\{q_L, d_L^c, u_L^c\}, \\ W &\longrightarrow e^{2i\beta}W\end{aligned}\tag{4}$$

so that ordinary quarks and Higgs doublets carry vanishing R charge.

3. R Violation and the μ Term

As a simple example, consider the Polonyi super-potential [13],

$$W_{Polonyi} = m^2(z + \tilde{m})\tag{5}$$

where \tilde{m} is of order of the Planck scale ($M = M_{Pl}/\sqrt{8\pi}$) and m is of the order of intermediate SUSY breaking scale (M_I). Supersymmetry is broken at the scale $m(\sim M_I)$, $F_z = m^2 + (m^2/M^2)z^*(z + \tilde{m})$ with $z = (\sqrt{3} - 1)M$ and $\tilde{m} = (2 - \sqrt{3})M$. With $\tilde{m} = 0$, $W = W_Y + W_{Polonyi}$ has the R symmetry with z transforming under R as

$$z \longrightarrow e^{2i\beta}z.\tag{6}$$

The field z gets mass through m of order $O(m^2/M) = O(m_{3/2})$. From the supergravity Lagrangian, we note [14] that the z field interacts with the other observable sector fields through gravitation only, and hence it has the severe cosmological problem [15]. The hypothetical R symmetry is broken by nonzero \tilde{m} . To see the effect of R symmetry violation, we can assign R -charge 2 to $m^2\tilde{m}$. Thus, we expect that this model generates through gravitational interaction a super-potential $(m^2\tilde{m})H_1H_2/M^2$, producing $\mu \sim m_{3/2}$. Any supergravity model, with super-potential $W = \text{cubic terms} + \text{constant}$, realizes this kind of R breaking and accordingly generates a nonzero μ term.

Let us briefly discuss the models discussed by Guidice and Masiero [8] and by Casas and Munoz [9] from this R symmetry argument. Guidice and Masiero suggest a Kähler potential

$$K = (H_1H_2 + h.c.) + \cdots\tag{7}$$

where \cdots denotes $\phi^*\phi$, etc. In supergravity models, the Guidice and Masiero solution gives an H_1H_2 term in the scalar potential, not to the superpotential. However, its effect can be absorbed in the superpotential through the transformation,

$$K(\phi, \bar{\phi}) \longrightarrow K(\phi, \bar{\phi}) - F(\phi) - \bar{F}(\bar{\phi}), \quad W(\phi) \longrightarrow e^{F(\phi)} W(\phi). \quad (8)$$

Therefore, the solutions discussed in Refs. [8,9] can be studied in the framework of Ref. [9]. Casas and Munoz assume that the superpotential W_0 does not contain the μ term. But the gravitational interaction can generate a term $W_0 H_1 H_2 / M^2$. They assume $\langle W_0 \rangle \sim M^2 m_{3/2}$ to generate $\mu \sim M_W$ and relate it to the scale of supersymmetry breakdown.³ The effective superpotential $W_0 H_1 H_2 / M^2$ preserves the R symmetry, and $\langle W_0 \rangle \neq 0$ breaks the R symmetry and there results a pseudo-Goldstone boson, which is hidden in the fields describing W_0 . This solution may have the Polonyi problem too [15].

A more interesting case is provided in models with hidden sector confining gauge groups [16].⁴ The μ term generated in this scenario has been discussed in Ref. [11]. We will elaborate this mechanism after presenting general classifications of R breaking mechanisms.

In string derived supergravity models, there exist only cubic terms in the superpotential. One can then define an R symmetry. This R symmetry could be broken⁵ in various ways: (a) vacuum expectation values of $R \neq 0$ scalar fields, (b) F terms of $R = 0$ chiral fields, (c) a constant in the effective superpotential, or (d) gaugino condensation.

Vacuum expectation values of $R \neq 0$ scalar fields – Let the chiral fields carrying nonzero R quantum numbers be B_i with $R = R_i$. Then the effective superpotential of the form

$$\frac{1}{M^{n-1}} [\prod_{i=1}^n B_i] H_1 H_2, \quad \sum_{i=1}^n R_i = 2 \quad (9)$$

³In this model there is no symmetry that forbids the μ -term and dangerous renormalizable operators like, for example, $Z_i H_1 H_2$ in W_0 with nonvanishing vacuum expectation values of gauge singlets Z_i . Thus, a full solution of the μ problem is not achieved.

⁴The most popular attempts of supersymmetry breaking in superstring models rely on this idea [17].

⁵In such theories a pseudo-Goldstone boson χ does appear. Its mass can be estimated by standard perturbative methods.

preserves the R symmetry. The R symmetry is broken by the vacuum expectation values of B_1, B_2, \dots , and B_n , and μ is generated

$$\mu = \frac{\langle B_1 B_2 \dots B_n \rangle}{M^{n-1}}. \quad (10)$$

For the argument of R symmetry toward a naturally small μ to make any sense, $\langle B_i \rangle \ll M$.

F terms of chiral fields – Let us suppose that A_i carry vanishing R quantum numbers so that A_{iF} carries $R = -2$. Then $A_i^* H_1 H_2$ needed for the D -term does not carry R quantum number and generates a μ term

$$\frac{1}{M} \int d^2 \bar{\theta} d^2 \theta A_i^* H_1 H_2 = \int d^2 \theta W_\mu \quad (11)$$

where $\mu = A_{iF}^*/M \sim m_{3/2}$. The dilaton superfields, moduli fields and chiral fields corresponding to flat directions can have vanishing R quantum numbers, since they do appear in the superpotential only as multiplicative factors of terms involving the interaction of the matter fields. For B_i fields carrying $R = \pm 1$, one needs two B 's, e.g. $B_1(R = 1)$ and $B_2(R = -1)$ to have a D -term,

$$\frac{1}{M^2} \int d^2 \bar{\theta} d^2 \theta B_1^* B_2^* H_1 H_2. \quad (12)$$

In this case,

$$\mu \sim (1/M^2)(B_{1F}^* \langle B_2^* \rangle + \langle B_1^* \rangle B_{2F}^*). \quad (13)$$

Since $\langle B_i \rangle \ll M$ for R symmetry to solve the μ problem, the μ term generated by B_i 's are insignificant compared to the one arising from the A_i terms.

Constants in W_{eff} – As discussed above the constant in the superpotential breaks the R symmetry. In string motivated supergravity models, this constant arises from the vacuum expectation values of $A_i B_j B_k$ where $R(A_i) = 0$ and $R(B_j, B_k) = 1$ or more generally $R(A_i) + R(B_j) + R(B_k) = 2$. For this to generate a gravitino mass scale μ , we require

$$\langle A_i \rangle \sim M, \quad \langle B_i \rangle \sim M_I. \quad (14)$$

Thus to obtain a reasonable μ , the scalar components of chiral fields participating in the cubic superpotential with nonzero R should have vacuum expectation values not exceeding the intermediate scale.

Gaugino condensation – Gauginos are assigned with $R = +1$ so that gauge bosons carry the vanishing R quantum number. Therefore, the gaugino condensation in the hidden sector breaks supersymmetry and can generate a μ term through

$$\frac{1}{M^2} \langle \Psi^a \Psi^a \rangle H_1 H_2 \quad (15)$$

where the (chiral) gauge boson multiplet is $\Psi^a = i\lambda^a + \dots$.⁶ The magnitude of μ is

$$\mu \sim \frac{\Lambda_h^3}{M^2}. \quad (16)$$

4. Solutions of the Strong CP and μ Problems in String Models

This leads us to the discussion on the anomalous $U(1)$ symmetry [19] in string models and generation of the μ term through the Peccei–Quinn symmetry [11]. The anomalous gauge $U(1)$ becomes a global $U(1)_P$ by removing the massive anomalous gauge boson [20].⁷ It obtains mass by absorbing the model-independent axion. Frequently, string models are referred to having no global symmetry, which led some to assume an approximate $U(1)_{PQ}$ in some string models [21]. However, the class of string models with an anomalous $U(1)_A$ have one $U(1)_P$ global symmetry below the string scale for a solution of the strong CP problem. The $U(1)_P$ symmetry has the P –(gauge boson)–(gauge boson) anomaly, [20,11],

$$\partial_\mu J_P^\mu = \frac{1}{32\pi^2} \sum_{G_i} F_{\mu\nu}(G_i) \tilde{F}^{\mu\nu}(G_i) \quad (17)$$

⁶One can obtain the same effect by generalizing the gauge kinetic function [18].

⁷The reason that the global symmetry results is similar to that an $SU(2)$ global symmetry results if an $SU(2)$ gauge symmetry without Yukawa couplings is broken by the vacuum expectation value of a doublet Higgs field.

where G_i is the gauge group. With one nonabelian gauge group [20], $U(1)_P$ is the needed Peccei–Quinn symmetry for the solution of the strong CP problem. But the hidden sector confining group needed for SUSY breaking invalidates the strong CP solution by the invisible axion since the axion in general gets a mass of order (coupling) $\times\Lambda_h^2$ where Λ_h is the scale of the hidden sector confining gauge group. In Ref. [11], it was pointed out that $U(1)_R$ can be used to obtain the invisible axion. We elaborate briefly how this idea can be made successful.

With $U(1)_P$ and $U(1)_R$ symmetries, one may consider a potential of the form

$$V = -\Lambda_{QCD}^4 \cos(\theta_c + \nu_0 \theta_h) - \Lambda_h^4 \cos(\theta_c + \nu_h \theta_h) + V_R(\theta_c, \theta_h) \quad (18)$$

where $\theta_c(= a_c/v_c)$ is the QCD vacuum angle and $\theta_h(= a_h/v_h)$ is the hidden sector vacuum angle. The Goldstone bosons corresponding to these symmetries are a_c and a_h with decay constants v_c and v_h , respectively. We assume that $U(1)_R$ has the $R - SU(3)_c - SU(3)_c$ and $R - G_h - G_h$ anomalies which are parametrized by ν_0 and ν_h , respectively. The $U(1)_P$ has both $U(1)_P - SU(3)_c - SU(3)_c$ and $U(1)_P - G_h - G_h$ anomalies. Without V_R , both θ_h and θ_c are settled to zero. One can expect the appearance of nontrivial V_R , and there exists a possibility that θ_c is not settled to zero. However, to study the extra potential V_R , one must consider the original $U(1)_P \times U(1)_R$ symmetry. Let us consider the following chiral and gauge fields as a simple example,

	A	B	Ψ^a	H_1	H_2	
P	2	-1	0	1	1	(19)
R	0	1	1	0	0	

where P and R are the charges of $U(1)_P$ and $U(1)_R$. The potential V_R arises from effective super-potential

$$\int d^2\theta W_{\text{eff}} = \int d^2\theta \{W_1 + \int d^2\bar{\theta} g_1 \bar{g}_2\} \quad (20)$$

where W_1 and g_1 are the composite superfield operator constructed from (left-handed) chiral fields only and \bar{g}_2 is constructed from (right-handed) anti-chiral fields only. From the $U(1)_P \times$

$U(1)_R$ symmetry, we argue that W_1 carries vanishing P and 2 units of R while $g_1\bar{g}_2$ carries vanishing P and R . Thus W_1 can contain ABB , BBH_1H_2/M , $\Psi^a\Psi^a$, etc. On the other hand $g\bar{g}$ can contain $A^*H_1H_2/M$, $A^*B^*B^*\Psi^a\Psi^a/M^3$, etc. From W_1 and $g_1\bar{g}_2$, we do not expect to generate a potential containing a_c and a_h since these phase fields do not appear in $\sum_i |\partial W_{\text{eff}}/\partial z_i|^2$ and $|W_{\text{eff}}|^2$. Thus $V_R(\theta_c, \theta_h) = (\text{constant})$. The constant is chosen so that the potential is zero at the minimum of the potential. Note, however, that V_R can contain the scalar partners of the phase fields.

To see the a_c and a_h independence of V_R , we observe that W_{eff} must preserve $U(1)_P \times U(1)_R$, which is the case for $U(1)_P$ if nonperturbative effects of gravitational interaction are supposed to respect it. Namely, we do not consider the argument based on wormholes, otherwise the axion solution of the strong CP problem is not attractive. Also, it is assumed that the $U(1)_R$ symmetry in some string models is valid up to dim=10 operators. This extra assumption is needed for the strong CP, but not for a solution of the μ problem. Suppose z_i carries P and Q charges so that it transforms under $U(1)_P$ and $U(1)_R$ as

$$z_i \longrightarrow e^{i[\alpha_c P(z_i) + \alpha_h R(z_i)]} z_i \quad (21)$$

where α_c and α_h are the rotation angles. The $U(1)_P$ and $U(1)_R$ symmetry of W_{eff} implies

$$P : W_{\text{eff}} \longrightarrow W_{\text{eff}} \quad (22)$$

$$R : W_{\text{eff}} \longrightarrow e^{i\alpha_h R} W_{\text{eff}} \quad (23)$$

where $R = \sum_i R(z_i) = 2$. The nonlinear realization of this transformation is represented by Goldstone fields a_c and a_h . W_{eff} does not depend on a_c but depends on a_h , which can be factored out as

$$W_{\text{eff}} = (W_{\text{eff}})_{\text{radial}} e^{2ia_h/v_h} \quad (24)$$

with the following transformation

$$a_h \longrightarrow a_h + \alpha_h v_h \quad (25)$$

where $(W_{\text{eff}})_{\text{radial}}$ does not contain a_h . $|W_{\text{eff}}|^2$ is independent of a_c and a_h . $\partial W_{\text{eff}}/\partial z_i$ carries $P = -P(z_i)$ and $R = 2 - R(z_i)$ and can be written

$$\left(\frac{\partial W_{\text{eff}}}{\partial z_i}\right)_{\text{radial}} \exp \left[i \left(-P(z_i) \frac{a_c}{v_c} + (2 - R(z_i)) \frac{a_h}{v_h} \right) \right] \quad (26)$$

where $(\partial W_{\text{eff}}/\partial z_i)_{\text{radial}}$ does not depend on a_c and a_h . Therefore, $\sum_i |\partial W_{\text{eff}}/\partial z_i|^2$ does not depend on a_c and a_h .

In general, we expect that ν_0 and ν_h are different, and both θ_h and θ_c can be settled to zero by the dynamical fields a_h and a_c . The mass matrix M^2 of a_c and a_h satisfies

$$\text{Det } M^2 = (\nu_0 - \nu_h)^2 \frac{\Lambda_{QCD}^4 \Lambda_h^4}{v_c^2 v_h^2} \quad (27)$$

from which we expect an invisible axion, with mass $\sim (\text{coupling constant}) \times \Lambda_{QCD}^2/\Lambda_h$, because $\nu_0 \neq \nu_h$ and Λ_h , v_c and v_h are of the same order. The other boson gets a mass of order $(\text{coupling constants}) \times \Lambda_h^2$.

The μ term given in Eq. (1) is supersymmetric. In principle, it can arise without supersymmetry breaking. Thus it can arise from a term AH_1H_2 in the superpotential with $\langle A \rangle \neq 0$, and μ can be of order string scale. In the above, we argued that a Peccei–Quinn symmetry spontaneously broken at the intermediate mass scale M_I excludes this possibility. Then there exist many possibilities for generating the μ term of order of M_{SUSY} .

Recently, the μ term has been calculated in some string models [22] indicating the presence of terms as described in the paper of Giudice and Masiero [8]. The contribution of terms in the superpotential to the μ term will, however, be dominant unless they are forbidden by a symmetry.

A general expression for the μ term can be written as

$$\mu = A + m_{3/2} O(1) - F^{\bar{n}} \partial_{\bar{n}} \xi(M^{\bar{n}}) + \tilde{\mu} \quad (28)$$

where F^n is the auxiliary component of n -th modulus, and $M^{\bar{n}}$ is the anti-modulus [22]. The A term must be absent, presumably by the Peccei–Quinn symmetry as mentioned above. The $m_{3/2}$ term arises from the Giudice–Masiero or the Casas–Munoz form [8,9]. The rest

arise from D terms; thus involve anti-chiral fields. ξ arises from moduli couplings and $\tilde{\mu}$ arises from the Yukawa and gauge couplings. In any case, these are at most of order $m_{3/2}$, since to obtain a superpotential W_μ one must take $\int d^2\bar{\theta}$ on the D -term which will pick up the F -term ($\sim M_I^2$) of anti-chiral fields.

Notice, that the absence of A in Eq. (28) is crucial for the solution of the μ problem and has to be guaranteed by a symmetry, as e.g. the anomalous $U(1)$ symmetry in the models discussed above.

5. Conclusion

We have seen that there can be many sources for the μ term in supergravity. To understand its magnitude in a natural scheme, however, one needs a symmetry. Both the R symmetry and the Peccei–Quinn symmetry are suggested for a natural solution of the μ problem. Without such a symmetry principle, one does not have a handle to remove specific (large) terms in the super-potential. We have shown that string models with the anomalous $U(1)_A$ with specific P and R charges are the best candidates toward the solution of both the strong CP problem and the μ problem.

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